

Spin Superfluidity in Biaxial Antiferromagnetic Insulators

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Antiferromagnets may exhibit spin superfluidity since the dipole interaction is weak. We seek to establish that this phenomenon occurs in insulators such as NiO, which is a good spin conductor according to previous studies. We investigate non-local spin transport in a planar antiferromagnetic insulator with a weak uniaxial anisotropy. The anisotropy hinders spin superfluidity by creating a substantial threshold that the current must overcome. Nevertheless, we show that the application of a high magnetic field removes this obstacle near the spin-flop transition of the antiferromagnet. Importantly, the spin superfluidity can then persist across many micrometers, even in dirty samples.

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Introduction.— Achieving long-range spin transport is essential in spintronics. In metals, conduction electrons can carry spin information. The spin-diffusion length is generally less than a few hundred nanometers and often as short as a couple of nanometers. However, in ferromagnets, there are additional transport channels via spin excitations, typically in the form of spin waves. In ferromagnetic insulators, the absence of noisy itinerant carriers implies less dissipation such that magnons can traverse distances up to several microns [1]. Magnetic low-damping insulators in which new spin transport mechanisms can exist are of interest and can be promising candidates in low-dissipation spintronics.

Antiferromagnets (AFMs) have ordered spin configurations, but there is no net magnetization at equilibrium. New observations and advances in our understanding have stimulated increased interest in AFM spintronics [2–4]. They produce no stray fields that can influence other elements. There are more known high-temperature AFM insulators and semiconductors than their ferromagnetic counterparts. AFMs exhibit transport properties similar to those of ferromagnets. Some of these features are anisotropic magnetoresistance [5], giant magnetoresistance [6], large anomalous Hall effect [7], and the spin Hall effect (SHE) [8]. There are also recent investigations of the spin Seebeck effect in AFMs [9–12]. Additionally, there are observations of spin transport in AFMs via spin pumping from an adjacent ferromagnet into AFMs [13–16]. In these experiments, it is possible that (evanescent) magnons carry the spin current [17]. A unique aspect of AFMs is that it is possible to trigger ultra-fast THz dynamics of the AFM order parameter via charge [18, 19] and spin currents [20], magnons [21], spin-orbit torques [22], light [23], and spin Hall oscillators [24]. Furthermore, magnon Bose-Einstein condensation (BEC) occurs in AFM insulators [25–27].

In this Letter, we investigate spin transport via spin superfluidity (SSF) in AFM insulators. Experiments have demonstrated that NiO is a good spin conductor [13, 15]. We therefore focus on NiO as a prototypical biaxial AFM insulator [28–30]. Crucially, the additional

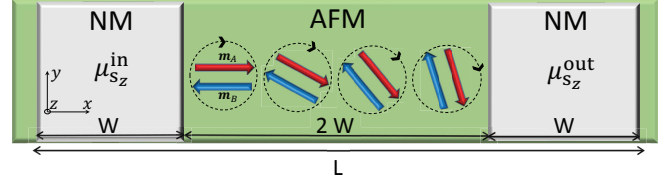


FIG. 1: SSF in a biaxial AFM insulator. The left and right normal metals (NMs) act as a spin injector and spin detector, respectively. The easy plane is the xy -plane and the easy axis is the x -direction. SSF occurs when the spins are tilted out of the easy plane and start to rotate around the hard axis with a spatially varying phase.

anisotropy normally hinders superfluidity [31–33]. However, we demonstrate that near the spin-flop transition [34] of biaxial AFMs, superfluid behavior nevertheless emerges. Importantly, we demonstrate through numerical calculations that long-range SSF beyond micrometers is feasible in real devices. This makes NiO and other biaxial AFMs promising for the first explicit experimental demonstration of SSF in magnetic materials.

Superfluidity is a dissipationless flow mediated by soft Nambu-Goldstone boson modes [35]. Models of superfluidity typically use a complex scalar field with global $U(1)$ symmetry. The superfluid velocity is proportional to the gradient of the condensate phase [35]. Halperin and Hohenberg demonstrated an analogy between the spin dynamics in planar magnetic systems and the hydrodynamic behavior of ideal liquids [36]. In a series of seminal works, Sonin extended the concept of superfluidity of electron-hole pairs [37, 38] to spin systems and introduced the notion of SSF [31]. In this scenario, SSF involves a 2π -rotation of spins in a planar magnet.

However, some of us have recently demonstrated that in planar ferromagnets, even as thin as 5 nm, the long-range dipole interaction destroys SSF based on the proposed mechanism [39]. Superfluidity reappears in synthetic AFM systems [39]. Since the dipole interaction is negligible in AFMs, we further explore SSF in AFMs [31, 40, 41].

The azimuthal angle of the order parameter, ϕ , and the out-of-plane component of the order parameter, n_z , can describe SSF [27, 31, 36, 39–41]. They are conjugate variables. In the superfluid phase, n_z is the superfluid density and ϕ is the phase of the condensate. The transverse component of the order parameter precesses. The spatial gradient of the superfluid phase is proportional to the spin supercurrent.

Setup.— We consider a quasi-one-dimensional (1D) biaxial AFM insulator. There is a strong hard-axis anisotropy and a weaker easy-axis anisotropy. Two metallic layers attach at the left and the right of the antiferromagnetic sample, as shown in Fig. (1). Inducing SSF requires the spin accumulation to be polarized along the hard anisotropy axis. The spin-valve structure proposed in Ref. [39] can achieve this requirement. It is also possible to use a heavy metal to create a spin current via SHE. In the latter case, there must be a finite angle between the hard axis and the interface normal to ensure a significant superfluid spin density.

We use a 1D Hamiltonian to model spin transport. The width of both leads is W , and the separation between them is $2W$. The total system length is $L = 4W$. A spin accumulation generated by one of the two mentioned injection methods induces a spin current from the left. In turn, the spin current exerts a torque on the spins in the AFM and causes precession. A small spatial gradient of the phase of the precession governs superfluidity. Finally, spin pumping into the right lead causes a spin accumulation therein. The resulting spin accumulation can be measured by either a spin-valve structure or the inverse spin Hall effect.

Magnetization dynamics and stability criteria.— We assume a collinear AFM with two equivalent magnetic sublattices. The unit vectors along the directions of the magnetic moments are $\mathbf{m}_A(\mathbf{r}, t)$ and $\mathbf{m}_B(\mathbf{r}, t)$, respectively. At equilibrium, $\mathbf{m}_A(\mathbf{r}, t)$ and $\mathbf{m}_B(\mathbf{r}, t)$ are antiparallel. Superfluidity can be described semiclassically. Then, the spin dynamics at each sublattice is described by the coupled Landau-Lifshitz-Gilbert (LLG) equations [2],

$$\dot{\mathbf{m}}_i = -\gamma \mathbf{m}_i \times \mathbf{H}_i^{\text{eff}} + \alpha(\mathbf{r}) \mathbf{m}_i \times \dot{\mathbf{m}}_i + \boldsymbol{\tau}_i^{\text{STT}}(\mathbf{r}), \quad (1)$$

where $i = A$ or B [42]. In Eq. (1), the first term on the right-hand side is the precessional torque driven by the effective magnetic field $\mathbf{H}_i^{\text{eff}}$, where γ is the effective gyromagnetic ratio. The second term, the damping torque, depends on the effective Gilbert coefficient $\alpha(\mathbf{r}) = \alpha_G + \alpha_{\text{SP}}(\mathbf{r})$. α_G is the dimensionless bulk Gilbert damping constant. $\alpha_{\text{SP}}(\mathbf{r})$ is the local damping enhancement arising from spin pumping into the normal metals [20, 45] within the two contact regions. The last term is the local spin transfer torque, $\boldsymbol{\tau}_i^{\text{STT}}(\mathbf{r}) = -\alpha_{\text{SP}}(\mathbf{r}) \mathbf{m}_i \times (\mathbf{m}_i \times \boldsymbol{\mu}_s^{\text{in}}(\mathbf{r}))$. $\boldsymbol{\mu}_s^{\text{in}}(\mathbf{r})$ is a spin accumulation at the left interface [16].

The free energy F determines the effective magnetic field, $\mathbf{H}_i^{\text{eff}} = -\delta F[\mathbf{m}_A, \mathbf{m}_B]/(\gamma \delta \mathbf{m}_i)$. In our system,

$F = \int d\mathbf{r} (f_{\text{ex}} + f_{\text{ani}} + f_Z)/V_0$, where V_0 is the AFM volume, resulting from the exchange interaction, the magnetic anisotropies, and the Zeeman energy:

$$f_{\text{ex}} = \omega_{\text{ex}} \mathbf{m}_A \cdot \mathbf{m}_B + \frac{\mathcal{A}^{(1)}}{2} [(\nabla \mathbf{m}_A)^2 + (\nabla \mathbf{m}_B)^2] + \mathcal{A}^{(2)} (\nabla \mathbf{m}_A \cdot \nabla \mathbf{m}_B), \quad (2a)$$

$$f_{\text{ani}} = \frac{\omega_{\perp}^{(1)}}{2} [(\mathbf{m}_A \cdot \hat{z})^2 + (\mathbf{m}_B \cdot \hat{z})^2] + \omega_{\perp}^{(2)} \mathbf{m}_A \cdot \hat{z} \mathbf{m}_B \cdot \hat{z} - \frac{\omega_{\parallel}^{(1)}}{2} [(\mathbf{m}_A \cdot \hat{x})^2 + (\mathbf{m}_B \cdot \hat{x})^2] - \omega_{\parallel}^{(2)} \mathbf{m}_A \cdot \hat{x} \mathbf{m}_B \cdot \hat{x}, \quad (2b)$$

$$f_Z = -\omega_H \mathbf{H} \cdot (\mathbf{m}_A + \mathbf{m}_B). \quad (2c)$$

In Eq. (2), ω_{ex} is the homogeneous exchange energy, $\mathcal{A}^{(1,2)}$ is the inhomogeneous exchange stiffness, $\omega_{\perp}^{(1,2)}$ is the hard-axis anisotropy energy, $\omega_{\parallel}^{(1,2)}$ is the uniaxial anisotropy energy, and ω_H is the Zeeman energy induced by an external magnetic field in an arbitrary direction.

We introduce the magnetization $\mathbf{m} = (\mathbf{m}_A + \mathbf{m}_B)/2$ and the staggered order parameter $\mathbf{n} = (\mathbf{m}_A - \mathbf{m}_B)/2$. Since $|\mathbf{m}_i| = 1$, $\mathbf{m} \cdot \mathbf{n} = 0$ and $\mathbf{m}^2 + \mathbf{n}^2 = 1$. We use the exchange approximation since the exchange energy is considerably larger than all other energy scales. In this limit, $|\mathbf{m}| \ll |\mathbf{n}|$. A standard calculation subsequently shows that the magnetization is a slave variable of the order parameter $\mathbf{m} = \dot{\mathbf{n}} \times \mathbf{n} / 4\omega_{\text{ex}} + \omega_H \mathbf{n} \times (\mathbf{H} \times \mathbf{n}) / 2\omega_{\text{ex}}$ [18, 19]. In terms of the staggered field, the total free energy density is [18, 19],

$$f = \lambda^2 \omega_{\parallel} (\nabla \mathbf{n})^2 + \omega_{\perp} (\mathbf{n} \cdot \hat{z})^2 - \omega_{\parallel} (\mathbf{n} \cdot \hat{x})^2 + \frac{\omega_H^2}{2\omega_{\text{ex}}} (\mathbf{H} \cdot \mathbf{n})^2 \quad (3)$$

where $\lambda = \sqrt{\mathcal{A}/\omega_{\parallel}}$ is the domain wall (DW) length, $\mathcal{A} = \mathcal{A}^{(1)} - \mathcal{A}^{(2)}$, $\omega_{\parallel} = \omega_{\parallel}^{(1)} - \omega_{\parallel}^{(2)}$, and $\omega_{\perp} = \omega_{\perp}^{(1)} - \omega_{\perp}^{(2)}$. The exchange stiffness \mathcal{A} and the homogeneous exchange energy ω_{ex} are related via $\omega_{\text{ex}}/\mathcal{A} = 2D/d^2$, where D and d are the spatial dimension and the lattice constant, respectively.

We consider an external magnetic field along the easy axis. Eq. (3) highlights that a critical magnetic field $H_c = \omega_H^c/\gamma = \sqrt{2\omega_{\text{ex}}\omega_{\parallel}}/\gamma$ compensates the uniaxial anisotropy. This peculiarity of AFM systems is known as the spin-flop transition [34]. At the spin-flop transition, the free energies of biaxial AFMs become similar to those of planar AFMs. However, the focus on this feature remains insufficient to prove, or fully understand, the range of validity of SSF.

The additional feature is that the magnetic field continues to influence the dynamic part of the AFM Lagrangian. This influence is via a gyrotropic term that breaks the Lorentz-invariant properties of AFMs, $\mathcal{L}_{\text{kin}}[\mathbf{n}] = (\dot{\mathbf{n}}^2 + 4\omega_H \dot{\mathbf{n}} \cdot \mathbf{n} \times \mathbf{H})/(8\omega_{\text{ex}})$. It is the total Lagrangian $\mathcal{L} = \mathcal{L}_{\text{kin}} - F$ that determines the dynamics of the Néel vector:

$$\mathbf{n} \times [\ddot{\mathbf{n}} - 8\lambda^2 \omega_{\text{ex}} \omega_{\parallel} \nabla^2 \mathbf{n} - 4\omega_H \mathbf{H} \times \dot{\mathbf{n}} + 8\omega_{\text{ex}} \omega_{\perp} n_z \hat{z} - 8\omega_{\text{ex}} \omega_{\parallel} n_x \hat{x} + 8\alpha \omega_{\text{ex}} \dot{\mathbf{n}} - 4\alpha_{\text{SP}} \omega_{\text{ex}} \mathbf{n} \times \boldsymbol{\mu}_s^{\text{in}}] = 0, \quad (4)$$

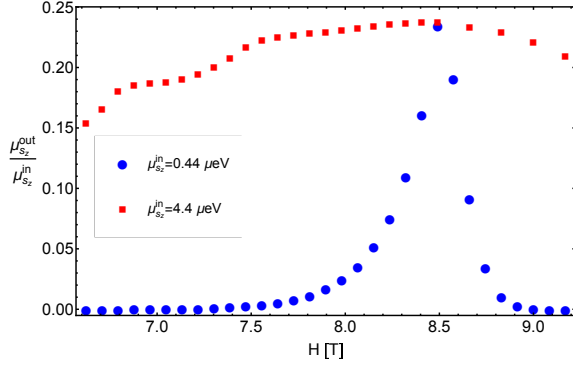


FIG. 2: The ratio between the output spin accumulation and the input spin accumulation as a function of the applied magnetic field for two different input spin accumulations.

where $\bar{\omega}_{\parallel} = \omega_{\parallel} - \omega_H^2/(2\omega_{\text{ex}})$ is the effective easy-axis anisotropy energy that has been renormalized by the magnetic field.

To study spin transport in the setup of Fig. (1), we consider one-dimensional solutions of the linearized equation of motion of Eq. (4). We use spherical coordinates for the staggered order parameter field $\mathbf{n} = \{\sqrt{1-n_z^2}\cos\phi, \sqrt{1-n_z^2}\sin\phi, n_z\}$.

First, we consider the static regime to find the critical current required to trigger SSF and the Landau criteria for the breakdown of SSF. To the linear order in n_z , we find

$$\bar{\omega}_{\parallel} \sin 2\phi - 2\lambda^2 \omega_{\parallel} \partial_x^2 \phi = -\alpha_{\text{SP}} \mu_z, \quad (5a)$$

$$(\omega_{\perp} + \bar{\omega}_{\parallel} \cos^2 \phi) n_z - \lambda^2 \omega_{\parallel} \partial_x^2 n_z = 0. \quad (5b)$$

where the driving force is $\mu_z = \mu_{s_z}^{\text{in}} \Theta(W-x)$ and Θ is the Heaviside function. One solution of Eq. (5) is a homogeneous state. The spins are then in the easy plane, $n_z = 0$, and the azimuthal angle is governed by the STT, $\sin 2\phi_0 = -(V_L/V_0)\alpha_{\text{SP}}\mu_{s_z}^{\text{in}}/\bar{\omega}_{\parallel}$, where V_L is the partial AFM volume below the left lead. The static macrospin solution becomes unstable when the spin-transfer torque is sufficiently large, $|\mu_{s_z}^{\text{in}}| > V_0\bar{\omega}_{\parallel}/(\alpha_{\text{SP}}V_L)$. Consequently, in the presence of a finite effective uniaxial anisotropy $\bar{\omega}_{\parallel}$, triggering SSF requires a large spin accumulation when the spin-pumping-enhanced damping α_{SP} is small. It is therefore essential to reduce the effective uniaxial anisotropy by an external magnetic field. In this regime, there are also two types of spatially varying solutions of Eq. (5). An in-plane homogeneous spiral state is a stable state when $\bar{\omega}_{\parallel} = 0$ while a kink-like state becomes more stable in $\bar{\omega}_{\parallel} \neq 0$ [33, 46].

Next, we allow the spins to vary in time. The dynamics, up to the linear order in n_z , is described by

$$\ddot{\phi} - v_c^2 (\partial_x^2 \phi - \frac{1}{2\lambda^2} \sin 2\phi) + 8\alpha_G \omega_{\text{ex}} \dot{\phi} + 4\omega_H \cos \phi \dot{n}_z = 0, \quad (6a)$$

$$\ddot{n}_z - v_c^2 \partial_x^2 n_z + 8\alpha_G \omega_{\text{ex}} \dot{n}_z + 8\omega_{\text{ex}} n_z (\omega_{\perp} + \bar{\omega}_{\parallel} \cos^2 \phi) - 4\omega_H \cos \phi \dot{\phi} = 0, \quad (6b)$$

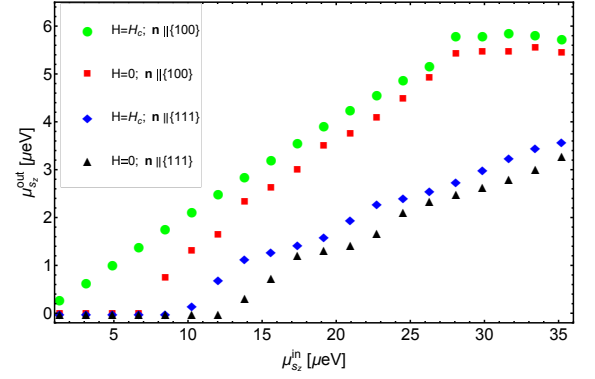


FIG. 3: The output spin accumulation (in the right lead) as a function of the input spin injection (in the left lead).

where $\bar{\lambda} = \sqrt{A/\bar{\omega}_{\parallel}}$ is the effective DW length, and the effective velocity of "light" is $v_c = 2\lambda\sqrt{2\omega_{\text{ex}}\omega_{\parallel}}$. The AFM spin wave velocity v_c is significantly larger than its ferromagnetic counterpart by a factor of $\sqrt{\omega_{\text{ex}}/\omega_{\perp}}$ [33].

Only after developing these equations (6a and 6b) can we explain why we can disregard the remaining magnetic field influence of the gyrotropic terms. The gyrotropic term in the Lagrangian causes the last term in Eq. (6a). This term couples the dynamics of the condensate phase to the superfluid density. However, this term is proportional to the superfluid density. We can therefore disregard this effect when a strong hard-axis anisotropy suppresses the out-of-plane dynamics.

Eqs. (6a) and (6b) decouple if $\omega_H = \omega_{\parallel} = 0$. Then, Eq. (6a) determines gapless phase excitations, $\omega_{\phi} = v_c k$. Conversely, Eq. (6b) implies that the density excitations have a gap, $\omega_{n_z} = \sqrt{v_c^2 k^2 + 8\omega_{\perp}\omega_{\text{ex}}}$. Since the gap is large, it is considerably more difficult to excite dynamic variations in the condensate density n_z . This is in sharp contrast to planar ferromagnets, where both phase and density excitations are gapless and have linear dispersion [40, 41, 43]. In the considered regime, $\omega_H = \omega_{\parallel} = 0$, the steady-state solution is $\phi(x, t) = \phi(x) + \Omega t$. The precession frequency is determined by the driving force of Eq. (5) such that $\alpha_G \Omega = \lambda^2 \omega_{\parallel} \partial_x^2 \phi$. The z -component of the magnetization is conserved. The continuity equation in the low Gilbert damping limit reads as $M_s \dot{n}_z = -\partial_x J_{s_z}$, where M_s is the saturation magnetization and the spin supercurrent density is $J_{s_z} = 2M_s \lambda^2 \omega_{\parallel} \partial_x \phi(x)$ [31, 40, 41]. As we have already discussed, this dependence of the spin superfluid current on the gradient of the condensate phase is similar to the characteristic of superfluid hydrodynamics of conventional superfluids.

In the presence of a finite effective uniaxial anisotropy, Eq. (6a) has a kink-like soliton solution for the condensate phase, $\phi(x, t) = 2 \tan^{-1} [\exp(\pm(x - x_0 - vt)/(\bar{\lambda}\sqrt{1 - (v/v_c)^2}))]$, where x_0 is the arbitrary DW center position and v is the DW velocity. An inhomogeneous state then becomes stable. The DW ve-

locity is determined by the driving force, resulting in $v \sim \pm W \alpha_{\text{SP}} \mu_s^{\text{in}} / (2\pi \alpha_G) < v_c$. The traveling solitons generate an AC signal on top of a DC output in the detector. At the spin-flop transition, $\bar{\lambda}^{-1} = 0$, only the DC component of the signal survives and the system exhibits perfect SSF.

Let us investigate the conditions for SSF stability. We will find two criteria that also exist in similar forms in ferromagnetic SSF [31]. To this end, we consider the limit of a large hard-axis anisotropy. The superfluid density n_z is then small, and the total free energy becomes $F \sim \lambda^2 \omega_{\parallel} (\partial_x \phi)^2 + [\omega_{\perp} - \lambda^2 \omega_{\parallel} (\partial_x \phi)^2 + \bar{\omega}_{\parallel} \cos^2 \phi] n_z^2 - \bar{\omega}_{\parallel} \cos^2 \phi$. At the spin-flop transition, $\bar{\omega}_{\parallel} = 0$, the free energy implies that the SSF remains stable provided that $\partial_x \phi < \sqrt{\omega_{\perp} / (\lambda^2 \omega_{\parallel})}$. This upper critical limit for the gradient of the condensate phase is analogous to the Landau criterion in conventional superfluidity. A finite effective anisotropy increases the upper critical limit. Another necessary condition is that the spin supercurrent must be spatially uniform. The steady-state solution of Eq. (6a) gives rise to an approximately uniform spin supercurrent only when $\partial_x \phi \gg 1/\bar{\lambda}$ [31].

Numerical results and discussion.— To establish SSF, we numerically solve the LLG equations for the setup in Fig. (1). In the numerical calculations, we use parameters for the prototypical AFM insulator NiO: $\omega_{\text{ex}}/2\pi \sim 27.4$ THz, $\omega_{\perp}/2\pi \sim 23$ GHz, $\omega_{\parallel}/2\pi \sim 1$ GHz, and $\alpha_G \sim 6.8 \times 10^{-3}$ [28–30]. With these parameters, the spin-flop transition occurs at a critical external magnetic field of $H_c \sim 8.5$ T, aligned along the easy-axis direction.

The spin accumulation pumped into the right lead, in the limit of low spin-memory loss, is $\mu_s^{\text{out}} = -(1/2) \sum_i \mathbf{m}_i \times \dot{\mathbf{m}}_i$ [45]. When the hard-axis anisotropy is large, the z -component of the pumped spin accumulation is given by $\mu_{s_z}^{\text{out}} \approx -\dot{\phi}$. In Fig. (2), we plot the normalized spin accumulation in the right lead versus the input spin accumulation in the left lead, $\mu_{s_z}^{\text{out}}/\mu_{s_z}^{\text{in}}$, as a function of the applied magnetic field along the x -direction H . The system length is $L = 0.75 \mu\text{m}$. We consider two STT amplitudes: $\mu_{s_z}^{\text{in}} = 0.44 \mu\text{eV}$ and $\mu_{s_z}^{\text{in}} = 4.4 \mu\text{eV}$. In both cases, the maximum spin-transport efficiency occurs at the spin-flop transition. When the STT energy is smaller than the easy-axis energy, the output signal vanishes rapidly when the applied magnetic field deviates from the critical field. Naturally, we still find large signals when the input STT energy is larger than the uniaxial energy. In the latter case, the STT is strong enough to overcome the pinning arising from the effective uniaxial anisotropy and triggers coherent spin dynamics even at lower external magnetic fields. This field-dependent behavior can be used to experimentally distinguish the spin current arising from usual magnons and the spin current carried by SSF.

In Fig. (3), we plot the detected spin accumulation as a function of the STT amplitude for a system size $L = 0.75 \mu\text{m}$. We show the results with and with-

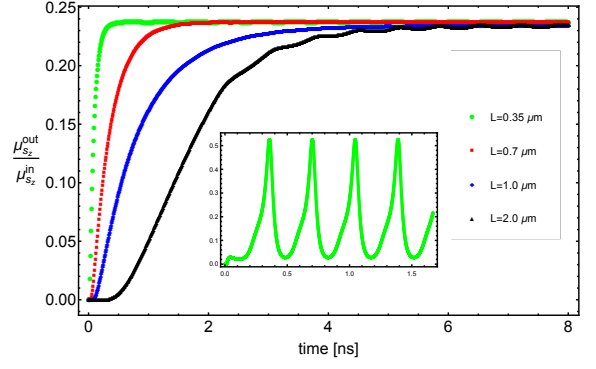


FIG. 4: The time evolution of the output spin accumulation for different length scales at the spin-flop transition $H = H_c$ when $\mu_{s_z}^{\text{in}} = 2 \mu\text{eV}$. Inset: Below the spin-flop transition $H \simeq 6.6$ T when $\mu_{s_z}^{\text{in}} = 4.4 \mu\text{eV}$.

out an applied critical magnetic field. When there is no magnetic field, the spins are pinned by the uniaxial anisotropy. In this case, the STT amplitude should be larger than a threshold that is proportional to the effective uniaxial anisotropy energy $\bar{\omega}_x$ to induce spin transfer. Near the spin-flop transition point, this limit vanishes due to the restoration of the U(1) symmetry. Above the threshold, the ratio between the pumped and injected spin currents is linear in the input spin accumulation, but only up to another critical value determined by the Landau criterion. There is no longer a typical superfluidity behavior beyond this point.

The magnetic moments in NiO are ferromagnetically aligned in the $\{111\}$ planes. The adjacent planes are antiferromagnetically coupled. Thus, the easy planes are not parallel to, e.g., a surface along $\{100\}$. We therefore also explore the likely effect of a finite angle between the easy plane of the NiO layer and a Pt/NiO interface. To this end, we rotate the easy plane in our numerical calculations while maintaining the applied magnetic field in the x -direction. Fig. (3) demonstrates that SSF remains feasible at finite angles. Only the component of the magnetic field parallel to the uniaxial anisotropy reduces the effective anisotropy. Then, in the presence of a finite angle between the external magnetic field and the uniaxial anisotropy θ , the critical magnetic field is increased to $H_c / \cos \theta$.

We plot the time evolution of the spin current for different length scales in Fig. (4). At the spin-flop transition point, even in a dirty sample where $\alpha_G \sim 6.8 \times 10^{-3}$ [28], SSF persists up to a few micrometers. This micron size range of the spin transport by SSF is considerably larger than the damping decay length of magnons in NiO, $\lambda_G = v_c / (4\alpha_G \omega_{\text{ex}}) \sim 40$ nm. Fig. (4) also shows that when the system size increases, the transient time increases. In general, in the presence of a uniaxial anisotropy and a strong STT amplitude, the output spin accumulation has both AC components in the GHz

regime together with a DC component. As discussed, the AC signal is a consequence of injecting kink-like solitons from the left lead. There is a reduction of the AC signal near the spin-flop transition. The signal is purely DC exactly at the transition point.

Experimentally, magnons and SSF contribute to the output spin accumulation. Their contributions can be distinguished either by changing the strength and direction of the magnetic field or by changing the sample size. Magnons decay exponentially with the system size, whereas we expect a very small algebraic decay for SSF in our setup.

Conclusion.— NiO is a biaxial AFM without a U(1) symmetry. This material can also have a significant Gilbert damping. These features appear to be detrimental to long-range spin transport via both magnons and SSF. Nevertheless, we demonstrate that NiO and other biaxial AFMs are good candidates for observing SSF over micrometer length scales. SSF behavior is dramatically improved around the spin-flop transition, which can be reached by applying an external magnetic field. SSF can be observed in standard non-local spin-transport setups and reach distances beyond micrometers.

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